

# Nonthermal Axino as Cool Dark Matter in Supersymmetric Standard Model without R-parity

Eung Jin Chun<sup>a\*</sup> and Hang Bae Kim<sup>b†</sup>

<sup>a</sup>*Korea Institute for Advanced Study*

*207-43 Cheongryangri-dong, Dongdaemun-gu, Seoul 130-012, Korea*

<sup>b</sup>*Department of Physics and Institute of Basic Science*

*Sungkyunkwan University, Suwon 440-746, Korea*

*\*ejchun@kias.re.kr, †hbkim@newton.skku.ac.kr*

## Abstract

We point out that the axino predicted as the supersymmetric partner of the axion is a good candidate for the recently proposed sterile neutrino cool dark matter. The axino mass falls into the right range ( $\lesssim$  keV) in the context of gauge mediated supersymmetry breaking. A sizable mixing of the axino with active neutrinos arises when R-parity violation is allowed and the resulting neutrino masses and mixing accommodate the atmospheric and solar neutrino oscillations simultaneously.

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In a recent paper [1], a new cosmological model has been proposed to explain the structure formation of the Universe. In this model, dark matter consists of non-thermal sterile neutrinos with masses about 100 eV to 10 keV which are produced around the big bang nucleosynthesis (BBN) epoch through a resonant active-sterile neutrino conversion in the presence of a net lepton number asymmetry. Such non-thermal neutrinos are called “cool” since their free streaming scale is larger than that of the usual cold dark matter but smaller than that of hot or warm dark matter. This cool dark matter (CtDM) model would be a good alternative to the conventional cold plus hot dark matter (CHDM) model which has been known to be the best cosmological model for the structure formation [2]. In view of particle physics theory, the validity of the CHDM model depends crucially on the origin of active neutrino masses and mixing. In the CHDM model, 20 % of dark matter is composed of hot components which are unarguably taken as three species of active neutrinos. The cosmological requirement,  $\Omega_\nu \simeq 0.2$ , constrains the sum of three active neutrino masses through the relation,  $\sum m_\nu = 92\Omega_\nu h^2$  eV, which implies  $\sum m_\nu \simeq 4.6$  eV, taking  $h = 0.5$ . If neutrinos take natural hierarchical masses as quarks and charged leptons, the recent atmospheric neutrino data from the Super-Kamiokande [3] imply  $\Delta m_{\text{atm}}^2 \approx m_{\nu_3}^2 < 10^{-2}$  eV<sup>2</sup> and thus  $\sum m_\nu < 0.1$  eV, in which case the CHDM model cannot work.

The purpose of this paper is to present a particle physics model in which the CtDM model is realized in a natural way while neutrinos obtain hierarchical masses which can accommodate the atmospheric [3] and solar neutrino data [4] simultaneously. One of popular ways to generate neutrino masses is to consider the minimal supersymmetric extension of standard model without imposing R-parity in which R-parity and lepton number violating bilinear and trilinear couplings give rise to typically hierarchical neutrino masses [5,6]. If one introduces in this context the Peccei-Quinn (PQ) mechanism which would be the most attractive solution to the strong CP problem [7], there exists a singlet fermion, called the axino  $\tilde{a}$ , which is a supersymmetric partner of the Goldstone boson, the axion, of a spontaneously broken PQ symmetry. Various cosmological and astrophysical observations are known to restrict the scale of PQ symmetry breaking:  $f_a \approx (10^{10} - 10^{12})$  GeV [8]. As was observed first in Ref. [9], the axino can be a sterile neutrino, that is, it mixes with active neutrinos once R-parity is not imposed. In the context of supergravity [10] where supersymmetry breaking is mediated at the Planck scale, the axino mass is quite model-dependent and could be in the range, 100 GeV – 1 keV [11]. However, in models with gauge mediated supersymmetry breaking (GMSB) [12] where the supersymmetry breaking mediation scale is much below the PQ scale  $f_a$ , the axino mass is determined by the Goldstone nature of the axion supermultiplet, and is predicted to be typically in the sub-keV regime in the minimal gauge mediation models [13].

We will show that the axino can be a good candidate of cool dark matter as its mass and mixing with active neutrinos fall into the right ranges for a reasonable range of parameter space in the context under consideration. In particular, the required mixing of the axino with active neutrinos can be obtained when one allows the R-parity violating terms whose sizes are fixed to explain the atmospheric and solar neutrino masses and mixing. A weak point of CtDM scenario would be the requirement of a large lepton asymmetry. It will be argued that the demanded lepton asymmetry can be generated through a late-time entropy production followed by the R-parity and lepton number violating decays of the ordinary lightest supersymmetric particle (LSP), or through the Affleck-Dine mechanism.

There are various sources for the cosmological axino population. The conventional source would be the thermal production. If the reheat temperature  $T_R$  after inflation is larger than the axino decoupling temperature  $T_d \sim f_a$ , thermally produced axinos may overclose the universe unless the axino mass is less than about 2 keV [14]. Therefore, the axino with mass  $\sim$  keV can be *warm* dark matter. When  $T_R < T_d$ , the primordial axinos are inflated away, but can be regenerated from thermal background. In this case, the axino *warm* dark matter needs the mass [15];

$$m_{\tilde{a}} \approx 3 \text{ MeV} \left( \frac{f_a}{10^{12} \text{ GeV}} \right)^2 \left( \frac{10^9 \text{ GeV}}{T_R} \right). \quad (1)$$

The other possibility is the secondary production of the axinos through the decay of ordinary superparticles. In a recent paper [16], it was argued that the axino with mass  $\sim 10$  GeV can be *cold* dark matter when the axinos are produced from the decay of the ordinary neutralino LSP. The last one would be the nonthermal production resulting in the axino *cool* dark matter, which we would like to discuss here.

Let us recapitulate the basic ingredients of the CDM model realizing the nonthermal production [1]. The production of a sterile neutrino  $\nu_s$  ( $= \tilde{a}$  in our scheme) which is much heavier than active neutrinos can be made by a resonant active-sterile neutrino conversion driven by a pre-existing large lepton asymmetry  $L = 10^{-4} - 10^{-1}$  which is destroyed during this process. The resonant oscillation between active and sterile neutrinos occurs at a temperature

$$T_{\text{res}} \approx 9 \left( \frac{m_{\nu_s}}{100 \text{ eV}} \right)^{1/2} \left( \frac{\mathcal{L}_i}{0.1} \right)^{-1/4} \left( \frac{E}{T} \right)^{-1/4} \text{ MeV}, \quad (2)$$

where  $\mathcal{L}_i = 2L_{\nu_i} + \sum_{j \neq i} L_{\nu_j}$ . The resonant transformation is adiabatic when the active-sterile neutrino vacuum mixing is not too small, that is,

$$\sin^2 2\theta_{si} \gtrsim 10^{-9} \quad (3)$$

for the lepton asymmetry  $L = 10^{-4} - 10^{-1}$ . Since the resonance temperature of a low energy neutrino is higher than that of a high energy neutrino, low energy neutrinos are produced at first consuming most of the lepton asymmetry. This makes high energy neutrino conversion non-adiabatic and no significant conversion occurs for high energy neutrinos. The produced sterile neutrinos are thereby cooler than the active neutrinos and the spectrum is non-thermal with  $\langle E \rangle / T \approx 0.7$ . The free-streaming length of these nonthermal sterile neutrinos is

$$\lambda_{\text{fs}} \sim \left( \frac{270 \text{ eV}}{m_{\nu_s}} \right) \left( \frac{\langle E \rangle / T}{0.7} \right) \text{ Mpc}, \quad (4)$$

and the contribution to the matter density today is given by

$$\Omega_{\nu_s} \approx \left( \frac{m_{\nu_s}}{343 \text{ eV}} \right) \left( \frac{0.5}{h} \right)^2 \left( \frac{\mathcal{L}_i}{0.1} \right). \quad (5)$$

In this mechanism, the sterile neutrino production can be confidently calculated only at the temperatures below the quark-hadron phase transition temperature, about 150 MeV. This translates to the condition

$$m_{\nu_s} < 23 \left( \frac{\mathcal{L}_i}{0.1} \right)^{1/2} \text{ keV}. \quad (6)$$

From Eqs. (5)–(6), one concludes that about 100 eV to 10 keV sterile neutrinos produced through the active–sterile neutrino conversion driven by a lepton asymmetry  $L \approx 10^{-4} - 10^{-1}$  can be a CDM candidate.

Our framework is the minimally extended supersymmetric standard model (MSSM) allowing R-parity and lepton number violating interactions which can explain the neutrino masses and mixing implied by the atmospheric and solar neutrino data. In order to estimate the axino mass and mixing with active neutrinos, we need to know the sizes of R-parity violating parameters determined from the neutrino data [3,4]. The R-parity and lepton number violating terms in the MSSM superpotential are

$$W = \epsilon_i \mu L_i H_2 + \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c, \quad (7)$$

where  $\mu$  is the Higgs mass parameter in the R-parity conserving superpotential,  $W = \mu H_1 H_2$ . As is well known [5], (active) neutrinos get masses at tree level due to the sneutrino vacuum expectation values (VEVs),  $\langle L_i^0 \rangle$ , misaligned with the bilinear terms  $\epsilon_i$ , as well as at 1-loop level through squark and slepton/Higgs exchanges. The tree-level neutrino mass matrix takes the form [17];

$$m_{ij}^{\text{tree}} \approx \xi_i \xi_j \frac{M_Z^2}{M_{1/2}} c_\beta^2, \quad (8)$$

where  $M_Z$  is the Z boson mass,  $M_{1/2}$  is the typical gaugino (photino or zino) mass, and  $c_\beta \equiv \cos \beta$ . Note that  $t_\beta \equiv \tan \beta = \langle H_2 \rangle / \langle H_1 \rangle$ . Here  $\xi_i \equiv \langle L_i^0 \rangle / \langle H_1 \rangle - \epsilon_i$  measure the misalignment between the sneutrino VEVs and  $\epsilon_i$ . Only one neutrino obtains a nonzero mass from the tree contribution (8) and the corresponding eigenvalue is given by  $m_{\nu_3} \approx \xi^2 (M_Z^2 / M_{1/2}) c_\beta^2$  where  $\xi \equiv \sqrt{\sum \xi_i^2}$ . Now, the atmospheric neutrino data from Super-Kamiokande [3] imply that  $m_{\nu_3} \approx \sqrt{\Delta m_{\text{atm}}^2} \approx 0.05$  eV and  $\xi_2 \approx \xi_3$  resulting in large mixing between the muon and tau neutrino. Therefore, one finds  $\xi \sim \xi_{2,3}$  and

$$\xi c_\beta \approx 10^{-6} \left( \frac{M_{1/2}}{M_Z} \right)^{1/2} \left( \frac{m_{\nu_3}}{0.05 \text{ eV}} \right)^{1/2}. \quad (9)$$

It is important for the later use to recall that nonzero values of  $\xi_i$  arise through renormalization group evolution from the mediation scale  $M_m$  to the weak scale and thus can be much smaller than the original parameters  $\epsilon_i$  or  $\lambda, \lambda'$  in Eq. (7). In one step approximation for integrating the renormalization group equation (See, for instances, Ref. [19].), one obtains

$$\xi_i \sim \epsilon_i \left( \frac{\mu A_b}{m_{\tilde{l}}^2} \right) \left( \frac{3h_b^2}{8\pi^2} \ln \frac{M_m}{m_{\tilde{l}}} \right) \quad (10)$$

where  $h_b, A_b$  are the bottom quark Yukawa coupling and the corresponding soft-parameter respectively, and  $m_{\tilde{l}}$  is the slepton mass. Taking  $\mu A_b = m_{\tilde{l}}^2$  and  $M_m = 10^3 m_{\tilde{l}}$ , one gets

$$\epsilon_i \sim 10^{-2} c_\beta \left( \frac{\xi_i c_\beta}{10^{-6}} \right). \quad (11)$$

The 1-loop contributions give nonzero masses to the other two neutrinos, which are much smaller than the tree contribution, in particular, in the context of GMSB models [18]. Among various 1-loop contributions to the second largest neutrino mass  $m_{\nu_2}$ , the diagrams with interchange of slepton and charged Higgs become important for large  $t_\beta$  and thus relevant for generating the solar neutrino mass scale. According to the estimation in Ref. [18], the mass ratio  $m_{\nu_2}/m_{\nu_3}$  is given by

$$\frac{m_{\nu_2}}{m_{\nu_3}} \sim 10^{-2} \left( \frac{\lambda_{233}}{\lambda'_{333}} \right) \left( \frac{t_\beta}{50} \right)^2 \quad (12)$$

under the assumption of the usual hierarchy among the trilinear couplings, that is, those involving third generations are larger than the others. The solar neutrino data explained by the matter resonant conversion, or the vacuum oscillation [4] requires  $m_{\nu_2} \approx \sqrt{\Delta m_{\text{sol}}^2} \sim 10^{-3}$ , or  $10^{-5}$  eV, respectively. Therefore, the matter conversion or vacuum oscillation solution can be obtained for  $t_\beta \sim 50$  or  $10$ , respectively, assuming  $\lambda_{233} \approx \lambda'_{333}$ . Note also that one needs  $\lambda_{233}, \lambda'_{233,333} \sim (10^{-4} - 10^{-6})(m_{\nu_2}/10^{-3} \text{ eV})^{1/2}$  for the right solar neutrino mass scale [6,18].

Let us now turn to the question why the axino can be CDM in the context of the R-parity violating MSSM with gauge mediated supersymmetry breaking. As discussed in Ref. [13], if a superfield  $S$  carries a PQ charge  $x_S$ , there arises the following effective Kähler potential for the superfield  $S$  and the axion superfield  $\Phi$ , which nonlinearly realizes the spontaneously broken PQ symmetry below its breaking scale  $f_a$ ,

$$K = \frac{x_S}{f_a} (\Phi^\dagger + \Phi) S^\dagger S. \quad (13)$$

Upon supersymmetry breaking, the interaction in Eq. (13) gives rise to a mass mixing between the axino and the fermionic component of  $S$  given by

$$m_D \approx x_S \frac{F_S}{f_a} \quad (14)$$

where  $F_S$  is the F-term of the field  $S$ . Without fine-tuning of some parameters, the value of  $F_S$  is expected to be of the order of  $M_S^2$  when the field  $S$  has the mass  $M_S$ . As a consequence, the axino gets a see-saw suppressed mass,  $m_{\tilde{a}} \approx m_D^2/M_S \sim x_S^2 M_S^3/f_a^2$ , that is,

$$m_{\tilde{a}} \sim 10 \text{ keV} \left( \frac{M_S}{10^5 \text{ GeV}} \right)^3 \left( \frac{10^{10} \text{ GeV}}{f_a} \right)^2. \quad (15)$$

This shows that the axino has a mass in the right range for being a CDM sterile neutrino, given that  $f_a = (10^{10} - 10^{12}) \text{ GeV}$  and  $M_S = (10^5 - 10^7) \text{ GeV}$  when  $S$  is a field in a messenger (or a hidden) sector of minimal gauge mediation models [20].

A sizable mixing of the axino with active neutrinos arises due to the bilinear terms in the superpotential (7) when the lepton doublets are charged under the PQ symmetry. Similarly to Eq. (14), the mixing mass between the axino and the active neutrino  $\nu_i$  is given by  $m_{si} \approx x_L F_L / f_a$  where  $x_L$  is the PQ charge of lepton doublet  $L$  and  $F_L \equiv \epsilon_i \mu \langle H_2 \rangle$ . Therefore, using the estimation of  $\epsilon_i$  in Eq. (11), one finds

$$\frac{m_{si}}{\text{keV}} \sim 0.1 c_\beta s_\beta \left( \frac{\mu}{500 \text{ GeV}} \right) \left( \frac{10^{10} \text{ GeV}}{f_a} \right) \left( \frac{\xi_i c_\beta}{10^{-6}} \right). \quad (16)$$

Given  $m_{\tilde{a}}$ , the adiabaticity condition  $\theta_{si} \approx m_{si}/m_{\tilde{a}} \gtrsim 10^{-5}$  (3) puts a lower bound on  $m_{si}$ , and an upper bound comes from the fact that the see-saw reduced mass of active neutrinos  $m_{si}^2/m_{\tilde{a}}$  should be smaller than the value  $m_{\nu_3} \sim 0.05 \text{ eV}$  used in Eq. (9). Taking roughly  $m_{si}^2/m_{\tilde{a}} < 10^{-2} \text{ eV}$ , we get the appropriate range of mixing mass

$$10^{-2} \left( \frac{m_{\tilde{a}}}{\text{keV}} \right) \text{ eV} \lesssim m_{si} \lesssim 3 \left( \frac{m_{\tilde{a}}}{\text{keV}} \right)^{1/2} \text{ eV}. \quad (17)$$

From Eqs. (15) and (16), one can find that the sterile and active neutrino mixing mass can be obtained in a reasonable range of the parameter space under consideration.

We have now to comment on the lifetime of our sterile neutrino, the axino, which should be stable as a dark matter. As we are working with gauge mediated supersymmetry breaking which has low supersymmetry breaking scale, the axino may be able to decay into a gravitino and an axion. For the supersymmetry breaking scale of order  $\sqrt{F} \lesssim 10^6 \text{ GeV}$  the gravitino mass is  $m_{3/2} \lesssim 1 \text{ keV}$  and thus the axino decay can be allowed kinematically. In this case, the lifetime of the axino is found to be [21]

$$\tau_{\tilde{a}} \approx 10^{29} \text{ sec} \left( \frac{1 \text{ keV}}{m_{\tilde{a}}} \right)^5 \left( \frac{\sqrt{F}}{10^5 \text{ GeV}} \right)^4. \quad (18)$$

That is, the axino is stable in the cosmological time scale for all the parameter ranges under consideration.

A nontrivial requirement for the CDM scenario is a large lepton asymmetry,  $L = 10^{-4} - 10^{-1}$ , around the nucleosynthesis period. This number is hierarchically larger than the baryon asymmetry  $B \simeq 10^{-10}$  required for producing the observed light element abundances through nucleosynthesis. Therefore, in the absence of a cancellation between  $L_{\nu_i}$ , the large lepton asymmetry must be generated at the temperature below the electroweak phase transition temperature to avoid the transfer between the lepton asymmetry and the baryon asymmetry due to the sphaleron effects. In our framework, since neutrino masses come from L violating couplings, they could also be the source of lepton asymmetry. Then the most promising way of generating such a lepton asymmetry would be the Dimopoulos-Hall mechanism [22].

When a late-time entropy production like thermal inflation is introduced for various cosmological reasons [23], the lepton asymmetry can be generated through the late-time decay of a (thermal) inflaton followed by lepton number violating superparticle decays [22]. The lepton asymmetry after the reheat of a late-time thermal inflation is then given by

$$L \simeq 5 \varepsilon_L \frac{T_R}{m_\phi} \quad (19)$$

where  $m_\phi$  is the mass of a thermal inflaton  $\phi$ ,  $T_R$  is the reheat temperature after the decay of  $\phi$ , and  $\varepsilon_L$  is the amount of lepton asymmetry per  $\phi$  decay. As a possibility to get a large  $\varepsilon_L$ , let us consider the case where the stau  $\tilde{\tau}$  is the ordinary LSP which occurs in a wide range of the GMSB parameter space [12]. As a stau is the LSP among the ordinary sparticles, it may decay into a tau and a gravitino ( $\tilde{\tau} \rightarrow \tau \tilde{G}$ ), or into a neutrino and a charged lepton ( $\tilde{\tau} \rightarrow \nu l$ ) through  $\lambda$  couplings in Eq. (7). As we want to have a large lepton asymmetry, the latter decay rate has to be larger than the former. In other words, we require  $\Gamma(\tilde{\tau} \rightarrow \nu l) \approx m_{\tilde{\tau}} |\lambda|^2 / 16\pi > \Gamma(\tilde{\tau} \rightarrow \tau \tilde{G}) \approx m_{\tilde{\tau}}^5 / 16\pi F^2$ , which puts a bound on  $\lambda$  depending on the supersymmetry breaking scale  $\sqrt{F}$  as follows:

$$\lambda > \frac{m_{\tilde{\tau}}^2}{F} \sim 10^{-6} \left( \frac{m_{\tilde{\tau}}}{100 \text{ GeV}} \right)^2 \left( \frac{10^5 \text{ GeV}}{\sqrt{F}} \right)^2. \quad (20)$$

In view of generating the solar neutrino mass scale  $m_{\nu_2}$  as given in Eq. (12) and the following discussions, the above inequality can be fulfilled for reasonable values of  $\sqrt{F}$ . The argument goes parallel with the stau decay into two quarks through  $\lambda'$  couplings which we neglect without loss of generality. As far as Eq. (20) is satisfied, the stau mainly decays through R-parity and lepton number violating interactions. Out-of-equilibrium condition can be met if the reheat temperature is smaller than  $m_{\tilde{\tau}}/20$  which is the typical supersymmetric particle decoupling temperature [8]. CP violation can come from complex phases of the leptonic charged current in case of Majorana neutrino masses, or from the complex L violating trilinear couplings  $\lambda, \lambda'$ . Let us now assume for simplicity that the right-handed stau  $\tilde{\tau}_R$  is the LSP which has decay modes  $\tilde{\tau}_R \rightarrow \nu_i e_j$ . Via the tree and one-loop interference depicted in FIG. 1, it gives

$$\varepsilon_L \approx \frac{2g_2^2 \text{Im}(\sum \lambda_{ij3} \lambda_{lm3}^* U_{il}^* U_{mj})}{4\pi \sum |\lambda_{ij3}|^2} \quad (21)$$

where  $U_{ij}$  is the CKM type mixing element for the charged leptonic current. The factor 2 comes from the fact that each  $\phi$  produces at least two staus. When for instance  $\lambda_{233}$  is the largest coupling, Eq. (21) becomes

$$\varepsilon_L \approx 2\alpha_2 \text{Im}(U_{22}^* U_{33}). \quad (22)$$

Note that  $U_{22}^* U_{33}$  is complex in general for Majorana neutrinos and its phase unconstrained by neutrino oscillation experiments can be of order one. Then the large mixing for atmospheric neutrino oscillations [3] implies  $\text{Im}(U_{22}^* U_{33}) \approx 1/2$ . Taking the upper limit of  $T_R \approx m_{\tilde{\tau}}/20$  which is required for a stau to be out of equilibrium, one finds

$$L \approx \frac{\alpha_2}{4} \frac{m_{\tilde{\tau}}}{m_\phi}. \quad (23)$$

Therefore, the maximal value of lepton asymmetry in this mechanism is  $L \sim 10^{-2}$ . For  $L \sim 10^{-3}$ , Eq. (6) implies  $m_{\tilde{a}} \lesssim 2 \text{ keV}$ . From Eq. (5), we obtain  $\Omega_{\tilde{a}} \approx 0.2$  for  $m_{\tilde{a}} \approx 2 \text{ keV}$  and  $L = 10^{-3}$ .

An even larger lepton asymmetry can arise through the Affleck-Dine mechanism [24] with a low reheat temperature. There is an interesting class of models where the lepton asymmetry  $L$  of order 1 can be produced through the Affleck-Dine mechanism. These are supergravity models that possess a so-called Heisenberg symmetry, which include no-scale type supergravity models and many string based supergravity models. In these models, supersymmetry breaking by the inflationary vacuum energy does not lift flat directions at tree level, and one-loop corrections gives small negative mass square for flat directions not involving large Yukawa couplings. After inflation these flat directions generate a large baryon and/or lepton asymmetry, typically of order 1, through the Affleck-Dine mechanism [25]. For example, we can use the well-known  $LH_2$  flat direction for this purpose, with the lepton number violation and CP violation arising from nonrenormalizable interactions. In this mechanism, it is possible to obtain  $L \approx 10^{-1}$  and hence  $\Omega_{\tilde{a}} \approx 1$ . One difficulty is that the large lepton asymmetry may accompany with the large baryon asymmetry. Several ways to reduce such large baryon asymmetry to the observed level were discussed in Ref. [26]. As we require the reheat temperature  $T_R$  to be below the electroweak scale, the decay produced, or the regenerated axino population can be completely negligible as can be seen from Eq. (1).

In conclusion, we have presented a well-motivated candidate for nonthermal sterile neutrino dark matter. The PQ mechanism realized in the supersymmetric standard model to resolve the strong CP problem predicts the presence of the axino, the fermionic superpartner of the axion. The lightness of the axino can arise when supersymmetry breaking is mediated by gauge interactions. We have pointed out that the axino mass can fall naturally into the demanded range  $(0.1 - 10)$  keV, given the range of the supersymmetry breaking scale,  $M_S \approx (10^4 - 10^7)$  GeV, and the PQ symmetry breaking scale,  $f_a \approx (10^{10} - 10^{12})$  GeV. It has also been shown that a sizable mixing between active neutrinos and the axino,  $\theta > 10^{-5}$ , required for it to be sterile neutrino dark matter, can arise from the (bilinear) R-parity violation which is introduced to explain the atmospheric and solar neutrino masses and mixing. An ‘ad hoc’ feature of nonthermal sterile neutrino dark matter scenario is that it needs a large lepton asymmetry  $L \approx 10^{-4} - 10^{-1}$ . We have argued that the lepton asymmetry up to  $L \sim 10^{-2}$  can be obtained by a late-time (thermal) inflaton decay followed by R-parity violating decays of the ordinary LSP  $\tilde{\tau}$ , given the maximal CP violating phases in the Majorana neutrino masses. An even larger lepton asymmetry can result from the Affleck-Dine mechanism. The reheat temperature is then required to be below the electroweak scale or a few GeV. The scenario under consideration needs better understanding why there is a big hierarchy between baryon and lepton asymmetries.



## REFERENCES

- [1] X. Shi and G. M. Fuller, Phys. Rev. Lett. **82**, 2832 (1999), astro-ph/9810076.
- [2] E. Gawise and J. Silk, Science **280**, 1405 (1998); M. A. K. Gross *et al.*, MNRAS **301**, 81 (1998).
- [3] The Super-Kamiokande Collaboration, Y. Fukuda, *et.al.*, Phys. Rev. Lett. **81**, 1562 (1998).
- [4] For a recent analysis, see J. N. Bahcall, P. I. Krastev and A. Yu. Smirnov, Phys. Rev. **D58**, 096016 (1998).
- [5] L. Hall and M. Suzuki, Nucl. Phys. **B231**, 149 (1984); A. S. Joshipura and M. Nowakowski, Phys. Rev. **D51**, 2421 (1995); F. M. Borzumati, Y. Grossman, E. Nardi and Y. Nir, Phys. Lett. **B384**, 123 (1996); R. Hempfling, Nucl. Phys. **B478**, 3 (1996); B. de Carlos and P. L. White, Phys. Rev. **D54**, 3427 (1996); A. Yu. Smirnov and F. Vissani, Nucl. Phys. **B460**, 37 (1996); M. Nowakowski and A. Pilaftsis, Nucl. Phys. **B461**, 19 (1996); H. P. Nilles and N. Polonsky, Nucl. Phys. **B484**, 33 (1997); E. Nardi, Phys. Rev. **D55**, 5772 (1997).
- [6] M. Drees, *et al.*, Phys. Rev. **D57**, 5335 (1998); R. Adhikari and G. Omanovic, Phys. Rev. **D59**, 073003 (1999); E. J. Chun, S. K. Kang, C. W. Kim and U. W. Lee, Nucl. Phys. **B544**, 89 (1999); V. Bednyakov, A. Faessler and S. Kovalenko, Phys. Lett. **B442**, 203 (1998); A. S. Joshipura and S. K. Vempati, hep-ph/9808232; Otto C.W. Kong, Mod. Phys. Lett. **A14**, 903 (1999); S. Rakshit, G. Bhattacharyya and A. Raychaudhuri, Phys. Rev. **D59**, 091701 (1999); D. E. Kaplan and Ann E. Nelson, hep-ph/9901254; Y. Nir and Y. Shadmi, JHEP 9905:023 (1999); A. Datta, B. Mukhopadhyaya and S. Roy, hep-ph/9905549.
- [7] J. E. Kim, Phys. Rep. **150**, 1 (1987).
- [8] E. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, 1990).
- [9] E. J. Chun, A. S. Joshipura and A. Y. Smirnov, Phys. Lett. **B357**, 608 (1995); Phys. Rev. **D54**, 4654 (1996).
- [10] H. P. Nilles, Phys. Rep. **110**, 1 (1984).
- [11] T. Goto and M. Yamaguchi, Phys. Lett. **B276**, 103 (1992); E. J. Chun, J. E. Kim and H. P. Nilles, Phys. Lett. **B287**, 123 (1992); E. J. Chun and A. Lukas, Phys. Lett. **B357**, 43 (1995).
- [12] For a review, see G. F. Giudice and R. Rattazzi, hep-ph/9801271.
- [13] E. J. Chun, hep-ph/9901220 (to appear in Phys. Lett. B).
- [14] K. Rajagopal, M. S. Turner and F. Wilczek, Nucl. Phys. **B358**, 447 (1991).
- [15] S. Chang and H. B. Kim, Phys. Rev. Lett. **77**, 591 (1996).
- [16] L. Covi, J. E. Kim and L. Roszkowski, Phys. Rev. Lett. **82**, 4180 (1999).
- [17] M. Nowakowski and A. Pilaftsis, in Ref. [5]; E. J. Chun and J. S. Lee, hep-ph/9811201 (to appear in Phys. Rev. D).
- [18] K. Choi, K. Hwang and E. J. Chun, hep-ph/9811363 (to appear in Phys. Rev. D); D. E. Kaplan and Ann E. Nelson, in Ref. [6].
- [19] E. Nardi, in Ref. [5]; A. S. Joshipura and S. K. Vempati, in Ref. [6].
- [20] M. Dine, A. Nelson, Y. Nir and Y. Shirman, Phys. Rev. **D53**, 1658 (1996).
- [21] E. J. Chun, H. B. Kim and J. E. Kim, Phys. Rev. Lett. **72**, 1956 (1994); H. B. Kim and J. E. Kim, Nucl. Phys. **B433**, 421 (1995).

- [22] S. Dimopoulos and L. Hall, Phys. Lett. **B196**, 135 (1987).
- [23] D. H. Lyth and E. D. Stewart, Phys. Rev. **D53**, 1784 (1996); E. D. Stewart, M. Kawasakia and T. Yanagida, Phys. Rev. **D54**, 6032 (1996); K. Choi, E. J. Chun and J. E. Kim, Phys. Lett. **B403**, 209 (1997); E. J. Chun, D. Comelli and D. H. Lyth, hep-ph/9903286.
- [24] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).
- [25] M. K. Gaillard, H. Murayama and K. A. Olive, Phys. Lett. **B355**, 71 (1995).
- [26] B. A. Campbell, M. K. Gaillard, H. Murayama and K. A. Olive, Nucl. Phys. **B538**, 351 (1999).

FIGURES

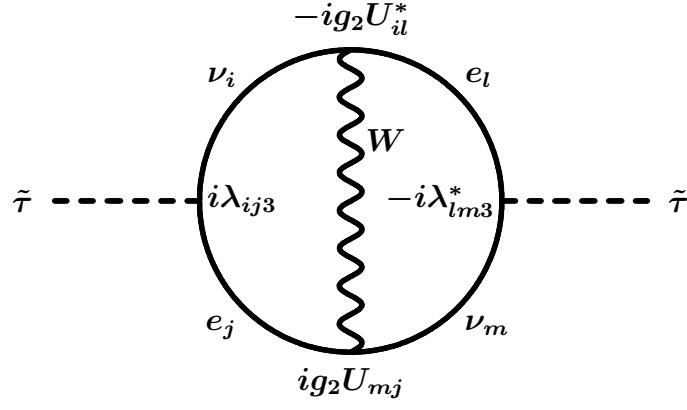


FIG. 1. The diagram for tree and one-loop interference